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INTRABAND RESONANCE SCATTERING OF ELECTROMAGNETIC RADIATION IN ANISOTROPIC QUANTUM DOTS

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We have developed a theory of the one-phonon intraband resonance scattering of electromagnetic radiation (IRSER) in anisotropic quantum dots subjected to an arbitrarily directed magnetic field. The differential cross section of scattering is obtained. The resonance structure of the cross section is studied. It is shown that the quantum dot subjected in a magnetic field can be used as the detector of phonon modes. The interesting multiplet structure of the resonance peaks is studied.

Keywords: quantum dot, resonance scattering.

1. Introduction

Theoretical [1–7] and experimental [8–12] studies of the resonance scattering of electromagnetic radiation in quantum dots (QD) taking into account the interaction with optical phonons are of great interest because the understanding of the scattering mechanisms is of fundamental importance for the applications. If the energy levels of the structure are discrete, the scattering becomes stronger resonantly. Such resonance scattering can provide the direct information on the electronic structure, phonon spectrum, and optical properties of QDs [13]. The most part of investigations is devoted to studying the *interband* resonance scattering. However, we suppose that the *intraband* resonance scattering is also of great interest because the distance between discrete levels in QDs can be done of order the optical phonon energy with help of the magnetic field. As a result we can use the magnetic field as the effective instrument to control optical and phonon properties QDs. It is important to note that the optical phonon emission is known to play a dominant role in QDs among the scattering mechanisms present in polar semiconductors.

Modern nanotechnology enables one to fabricate quantum dots of different shapes. In the past years the significant interest has been given to quantum wells and QDs that are characterized by an asymmetric confining potential [14–16]. In this work we present a theoretical study of the intraband resonance scattering of electromagnetic radiation in an anisotropic quantum dot subjected to a uniform magnetic field arbitrarily directed with respect to the potential symmetry axes. The applied magnetic field gives us the possibility to change the distance between levels and to adjust the energy levels of QDs on the various phonon modes. The study of the different polarization for the incident and emitted radiation yields the additional information about the phonon spectrum. Note that the study of IRSER lets us to obtain the simple analytic formulae for the cross-section in the case of anisotropic QDs.

IRSER in our case can be qualitatively described in the following way: the absorption of quantum $\hbar\omega_i$ of the high-frequency field (laser pump), emission of optical phonon $\hbar\omega_{\mathbf{q}}$ (photon $\hbar\omega_s$) in an intermediate state and emission of photon $\hbar\omega_s$ (optical phonon $\hbar\omega_{\mathbf{q}}$) in the initial state (see Fig. 1). In this approach, the cross section of resonant scattering is calculated from third order time-dependent perturbation theory.

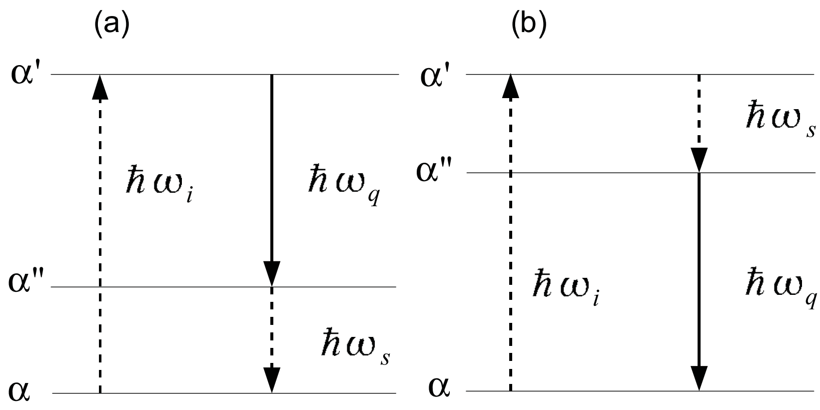

FIG. 1. Transitions leading to resonant absorption

We model the semiconductor QD using the asymmetric parabolic confinement $V(r) = m^*(\Omega_x^2 x^2 + \Omega_y^2 y^2 + \Omega_z^2 z^2)/2$ (here m^* is the electron effective mass, Ω_i ($i = x, y, z$) are the characteristic frequencies of the parabolic potential). The spectrum of electrons in such dot placed in an arbitrarily directed magnetic field $\mathbf{B}$ with the vector potential $\mathbf{A} = (B_y z/2 - B_z y, 0, B_x y - B_y x/2)$ has the form $\varepsilon_{nml} = \hbar\omega_1(n + 1/2) + \hbar\omega_2(m + 1/2) + \hbar\omega_3(l + 1/2)$ ($n, m, l = 0, 1, 2, \ldots$), where hybrid frequencies ω_j ($j = 1, 2, 3$) are obtained from the sixth-order algebraic equation [18].

2. Differential cross section

The differential resonance cross section $d^2\sigma/d\Omega d\omega_s$ of a volume V per unit solid angle $d\Omega$ for incident radiation with the frequency ω_i and emitted radiation with the frequency ω_s is given by [5] in analogy with the Raman cross section

$$\frac{d^2\sigma}{d\Omega d\omega_s} = \frac{V^2\omega_s^3 n_i n_s^3}{8\pi^3 c^4 \omega_i} W(\omega_s, \mathbf{e}_s) \tag{1}$$

where $n_i(n_s)$ is the refractive index of the medium with frequency ω_i (ω_s), c is the velocity of light, $\mathbf{e}_s$ is the unit polarization vector of the emitted radiation. The transition rate is calculated according to

$$W(\omega_s, \mathbf{e}_s) = \frac{2\pi}{\hbar} \sum_\alpha |W_{\alpha\alpha}|^2 \,\delta(\hbar\omega_i - \hbar\omega_s - \hbar\omega_q), \tag{2}$$

where $\alpha = (n, m, l)$ are the quantum numbers of the initial electron states in QD.

We consider only resonance transitions. In this case the scattering amplitude probability for phonon emission in QDs is described by a sum of two terms

$$\begin{aligned} W_{\alpha\alpha} = {} & \sum_{\alpha',\alpha''} \frac{\langle\alpha|\hat{H}_R(\omega_s)|\alpha''\rangle\langle\alpha''|\hat{H}_L|\alpha'\rangle\langle\alpha'|\hat{H}_R(\omega_i)|\alpha\rangle}{(\varepsilon_{\alpha'} - \varepsilon_\alpha - \hbar\omega_i)(\varepsilon_{\alpha''} - \varepsilon_\alpha - \hbar\omega_i + \hbar\omega_{\mathbf{q}})} \\ & + \sum_{\alpha',\alpha''} \frac{\langle\alpha|\hat{H}_L|\alpha''\rangle\langle\alpha''|\hat{H}_R(\omega_s)|\alpha'\rangle\langle\alpha'|\hat{H}_R(\omega_i)|\alpha\rangle}{(\varepsilon_{\alpha'} - \varepsilon_\alpha - \hbar\omega_i)(\varepsilon_{\alpha''} - \varepsilon_\alpha - \hbar\omega_i + \hbar\omega_s)} \end{aligned} \tag{3}$$

The first term in Eq. (3) corresponds to the transitions depicted on Fig. 1a, the second term in Eq. (3) corresponds to the transitions depicted on Figure 1b.

Here H_L is the operator electron-phonon interaction

$$\hat{H}_L = \sum_{\mathbf{q}} D_{\mathbf{q}} C_{\mathbf{q}} \exp(i\mathbf{q}\mathbf{r}) + \text{c.c}, \tag{4}$$

where $D_{\mathbf{q}}$ is the electron-phonon coupling constant.

The operator of the electron-photon interaction can be expressed as

$$\hat{H}_R = \frac{e}{m^*}\sqrt{\frac{2\pi\hbar N}{\varepsilon\omega}}\mathbf{e}\mathbf{P}, \tag{5}$$

where N is the number of initial-state photons with frequency ω, $\mathbf{e}$ is the polarization vector and $\mathbf{P} = \mathbf{p} - e\mathbf{A}/c$ is the generalized momentum, ε is the real part of the dielectric constant.

A direct calculation of the matrix elements of electron-photon and electron-phonon interactions is a complicate problem in our case. However, the method of canonic transformation of the phase space allows us to resolve this problem using only simple calculations from linear algebra [18]. In particular, in our preceding papers we used this method to study hybrid [19], hybrid-phonon [20] and hybrid-impurity resonances in this system [21].

Using the results obtained in [19], we can write the matrix elements of the operator $\hat{H}_R$ in the following form

$$\begin{gathered}\langle n'm'l'|\hat{H}_R|nml\rangle = ie\hbar\sqrt{\frac{\pi N}{m^*\varepsilon\omega}} \times \Big[X_1\sqrt{n+1}\delta_{n',n+1}\delta_{m',m}\delta_{l',l} \\ +X_2\sqrt{m+1}\delta_{n,n'}\delta_{m',m+1}\delta_{l',l} + X_3\sqrt{l+1}\delta_{n',n}\delta_{m',m}\delta_{l',l+1}\Big].\end{gathered} \tag{6}$$

where the coefficients X_i $(i = 1, 2, 3)$ were found in [19].

We introduce the notation

$$\begin{gathered}J(n''m''l'', n'm'l') = \left(\frac{n''!m''!l''!}{n'!m'!l'!}\right)^{1/2}(-1)^{n'-n''} \\ \times(-1)^{m'-m''}(-1)^{l'-l''}\exp[i\varphi_1(n'-n'')] \\ \times\exp[i\varphi_2(m'-m'')]\exp[i\varphi_3(l'-l'')]g_1^{n'-n''}g_2^{m'-m''} \\ \times g_3^{l'-l''}L_{n''}^{n'-n''}(g_1^2)L_{m''}^{m'-m''}(g_2^2)L_{l''}^{l'-l''}(g_3^2).\end{gathered} \tag{7}$$

Here $g_j = \sqrt{\lambda_j^2 + \kappa_j^2 l_j^4}/\sqrt{2}l_j$, $\tan\varphi_j = \kappa_j l_j^2/\lambda_j$, $l_j = \sqrt{\hbar/m^*\omega_j}$ $(j = 1, 2, 3)$ are the hybrid lengths, , $L_n^{n'}(x)$ are the generalized Laguerre polynomials, $\lambda_i = \hbar(b_{1i}q_x + b_{2i}q_y + b_{3i}q_z)$ $(i = 1, 2, 3)$, $\kappa_{i-3} = b_{1i}q_x + b_{2i}q_y + b_{3i}q_z$ $(i = 4, 5, 6)$, b_{ji} are components of the transition matrix [21].

Using (7), we can write the matrix elements of the operator electron-phonon interaction as

$$\begin{gathered}\langle n'm'l'|\hat{H}_L|n''m''l''\rangle = \sum_{\mathbf{q}} D_q\sqrt{N_{\mathbf{q}}}\exp(-g^2/2) \\ \times\exp[-(\kappa_1\lambda_1 + \kappa_2\lambda_2 + \kappa_3\lambda_3)i/2]J(n''m''l'', n'm'l'),\end{gathered} \tag{8}$$

where $N_{\mathbf{q}}$ is the number of phonons with the wave vector $\mathbf{q}$ and $g^2 = g_1^2 + g_2^2 + g_3^2$.

Substituting (6) and (8) into (3) after some cumbersome algebra it is possible to get analytic expression for $W_{\alpha\alpha}$. We consider only the resonance scattering. In this case the frequency of the pump is equal to the distance between the levels of QD. For definiteness, assume that we pump the QD by the laser with frequency $\omega_i = \omega_1$. Then we need to keep only resonance terms in the formula for $W_{\alpha\alpha}$. In this case for the first term in Eq. (3), we obtain

$$\begin{aligned}
&\sum_{\alpha',\alpha''} \frac{\langle\alpha|\hat{H}_R(\omega_s)|\alpha''\rangle\langle\alpha''|\hat{H}_L|\alpha'\rangle\langle\alpha'|\hat{H}_R(\omega_i)|\alpha\rangle}{(\varepsilon_{\alpha'} - \varepsilon_\alpha - \hbar\omega_i)(\varepsilon_{\alpha''} - \varepsilon_\alpha - \hbar\omega_i + \hbar\omega_{\mathbf{q}})} \\
&= -\frac{\pi e^2}{m^*\varepsilon}\sqrt{\frac{N_i(N_s+1)}{\omega_i\omega_s}} \sum_{\mathbf{q}} D_q \sqrt{N_{\mathbf{q}}} \exp(-g^2/2) \\
&\times \exp[-(\kappa_1\lambda_1 + \kappa_2\lambda_2 + \kappa_3\lambda_3)i/2] \frac{\sqrt{n+1}X_1^i}{\omega_1 - \omega_i} \\
&\times \left[\frac{\sqrt{m+1}X_2^S J(nm+1l, n+1ml)}{\omega_2 - \omega_i + \omega_{\mathbf{q}}} + \frac{\sqrt{l+1}X_3^S J(nml+1, n+1ml)}{\omega_3 - \omega_i + \omega_{\mathbf{q}}}\right].
\end{aligned} \tag{9}$$

Here the indexes i and s are referred to the incident and emitted photons, respectively.

The second term in (3) has the following form

$$\begin{aligned}
&\sum_{\alpha',\alpha''} \frac{\langle\alpha|\hat{H}_L|\alpha''\rangle\langle\alpha''|\hat{H}_R(\omega_s)|\alpha'\rangle\langle\alpha'|\hat{H}_R(\omega_i)|\alpha\rangle}{(\varepsilon_{\alpha'} - \varepsilon_\alpha - \hbar\omega_i)(\varepsilon_{\alpha''} - \varepsilon_\alpha - \hbar\omega_i + \hbar\omega_s)} \\
&= -\frac{\pi e^2}{m^*\varepsilon}\sqrt{\frac{N_i(N_s+1)}{\omega_i\omega_s}} \sum_{\mathbf{q}} D_q \sqrt{N_{\mathbf{q}}} \exp(-g^2/2) \\
&\times \exp[-(\kappa_1\lambda_1 + \kappa_2\lambda_2 + \kappa_3\lambda_3)i/2] \frac{\sqrt{n+1}X_1^i}{\omega_1 - \omega_i} \\
&\times \left[\frac{\sqrt{m}X_2^S J(nml, n+1m-1l)}{\omega_1 - \omega_2 - \omega_i + \omega_s} + \frac{\sqrt{l}X_3^S J(nml, n+1ml-1)}{\omega_1 - \omega_3 - \omega_i + \omega_s}\right].
\end{aligned} \tag{10}$$

Now we need to sum these terms to get $W_{\alpha\alpha}$. Taking into account the conversation law $\hbar\omega_i = \hbar\omega_{\mathbf{q}} + \hbar\omega_s$ we can transform denominators in Eq. (9) and Eq. (10). In this case we get for $W_{\alpha\alpha}$ the following formula

$$\begin{aligned}
W_{\alpha\alpha} &= -\frac{\pi e^2}{m^*\varepsilon}\sqrt{\frac{N_i(N_s+1)}{\omega_i\omega_s}} \sum_{\mathbf{q}} D_q \sqrt{N_{\mathbf{q}}} \exp(-g^2/2) \\
&\times \exp[-(\kappa_1\lambda_1 + \kappa_2\lambda_2 + \kappa_3\lambda_3)i/2] \frac{\sqrt{n+1}X_1^i}{\omega_1 - \omega_i} \\
&\times \left\{\frac{X_2^S}{\omega_2 - \omega_s}\left[\sqrt{m+1}J(nm+1l, n+1ml) - \sqrt{m}J(nml, n+1m-1l)\right]\right. \\
&+ \left.\frac{X_3^S}{\omega_3 - \omega_s}\left[\sqrt{l+1}J(nml+1, n+1ml) - \sqrt{l}J(nml, n+1ml-1)\right]\right\}
\end{aligned} \tag{11}$$

We can transform the differences in Eq. (11) using the recurrent formula for the Laguerre polynomials $xL_n^{\alpha+1} = (n+\alpha+1)L_n^\alpha(x) - (n+1)L_{n+1}^\alpha(x)$. As a result we get for the first difference

$$\begin{aligned}
&\sqrt{m+1}J(nm+1l, n+1ml) - \sqrt{m}J(nml, n+1m-1l) \\
&= -\exp(i\varphi_1)\exp(-i\varphi_2)g_1g_2L_n^1(g_1^2)L_m(g_2^2)L_l(g_3^2)
\end{aligned} \tag{12}$$

The second difference in Eq. (11) is calculated in analogy with Eq. (12). Using the obtained estimation, we get the following formula for the square of the scattering amplitude probability

$$\begin{aligned}|W_{\alpha\alpha}|^2 = \frac{\pi^2 e^4}{m^{*2}} \frac{N_i(N_s+1)}{\omega_i \omega_s} \sum_{\mathbf{q}} |D_{\mathbf{q}}|^2 N_{\mathbf{q}} \exp(-g^2) g_1^2 \\ \times [L_n^1(g_1^2)]^2 [L_m(g_2^2)]^2 [L_l(g_3^2)]^2 \frac{(n+1)(X_1^i)^2}{(\omega_1-\omega_i)^2} \\ \times \left| \frac{g_2 X_2^S}{\omega_2-\omega_s} \exp(-i\varphi_2) + \frac{g_3 X_3^S}{\omega_3-\omega_s} \exp(-i\varphi_3) \right|^2 .\end{aligned} \tag{13}$$

Then we can write the final formula for the cross-section taking into account the smearing of the hybrid levels by collisions

$$\begin{aligned}\frac{d^2\sigma}{d\Omega d\omega_s} = \frac{V\omega_s^2 n_i e^4 N_i(N_s+1)}{4\hbar^2 n_s c^4 m^{*2}} \sum_{\mathbf{q}} |D_{\mathbf{q}}|^2 N_{\mathbf{q}} \exp(-g^2) \\ \times g_1^2 [L_n^1(g_1^2)]^2 [L_m(g_2^2)]^2 [L_l(g_3^2)]^2 \frac{(n+1)(X_1^i)^2}{(\omega_1-\omega_i)^2+\Gamma^2} \\ \times \left| \frac{g_2 X_2^S}{\omega_2-\omega_s-i\Gamma} \exp(-i\varphi_2) + \frac{g_3 X_3^S}{\omega_3-\omega_s-i\Gamma} \exp(-i\varphi_3) \right|^2 \\ \times \delta(\omega_i-\omega_s-\omega_q),\end{aligned} \tag{14}$$

where Γ is the lifetime broadening.

3. Results and discussions

Equation (14) clearly shows that if one ignores the optical phonons dispersion and if the frequency of the pump is equal to ω_1 then we have the input resonance on the frequency ω_1 and the output resonance on the frequencies ω_2 and ω_3. Note that it is forbidden transitions with simultaneous changing more than one quantum numbers in the case of absorption (emission) of photon due to the selection rules. The possible transitions is shown on the Figure 2.

It is important to note that the hybrid frequency ω_k ($k = 1, 2, 3$) is determined by the magnitude and the direction of the magnetic field (i.e. they can be tuned with the help of the magnetic field). Hence, using the tunable laser and changing, for example, the magnitude of the magnetic field we can register phonon modes (with frequencies $\omega_q = \omega_1 - \omega_2$ and $\omega_q = \omega_1 - \omega_3$) in quantum dots as series resonance peaks in the dependence of the cross section on the magnetic field. The frequency of the phonon mode can be determined from the dependence of the magnetic field on the hybrid frequencies.

Let us now to study effects arising due to the dispersion of the phonons. Replacing the sum over the phonon wave vector by the integral in Equation (14) and assuming a parabolic dispersion low for long-wave phonons $\omega_q = \omega_0(1 - \omega_0^{-2} V_s^2 q^2)$, where ω_0 is the optical-phonon threshold frequency and V_s is the speed of sound, we can easily evaluate the integral with respect to $|\mathbf{q}|$ thanks to the presence of a delta function $\delta(\omega_i - \omega_s - \omega_q)$.

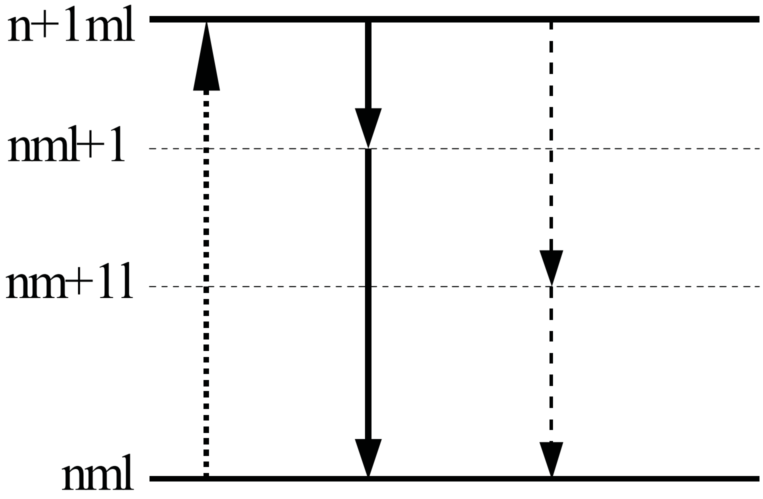

FIG. 2. Possible transitions leading to the resonant absorption in the case of anisotropic quantum dots. The dotted curve corresponds to the absorption of the pump field with the frequency ω_1. The solid line corresponds to the transitions with the change of the quantum number m. The dashed curve corresponds to the transitions with the change of the quantum number l.

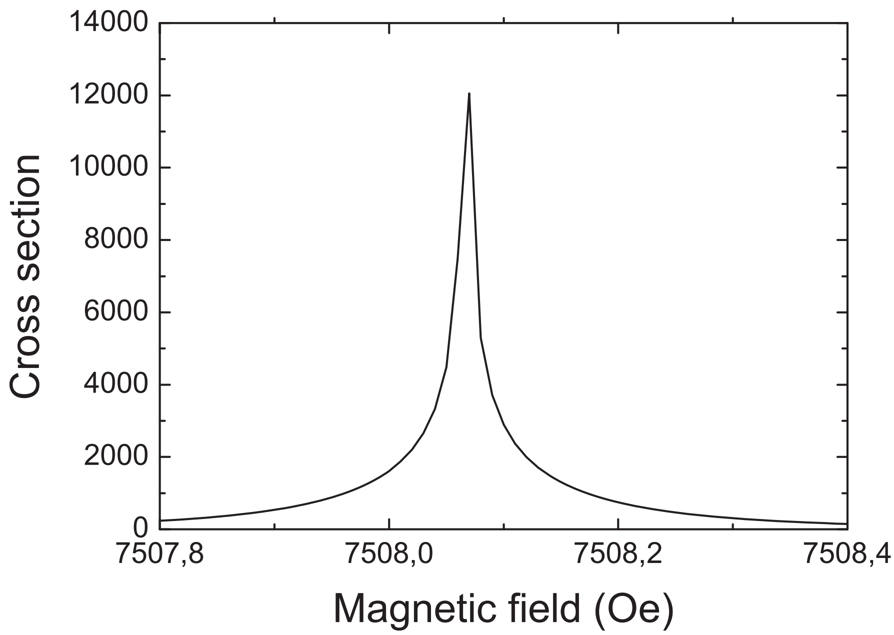

FIG. 3. Differential cross section (in arb. units) as a function of a magnetic field in the case of transition from the ground state and emission of PO-phonons, $\omega_0 = 1.2\times10^{12}$ s^{-1}, $\omega_x = 1.3\times10^{12}$ s^{-1}, $\omega_y = 1.4\times10^{12}$ s^{-1}, $\omega_z = 7.1\times10^{13}$ s^{-1}.

Converting to spherical coordinates we obtain the following equation for the differential cross section

$$\begin{aligned}\frac{d^2\sigma}{d\Omega d\omega_s} &= \frac{V\omega_s^2 n_i e^4 N_i(N_s+1)N_0\omega_0^{3/2}}{8\hbar^2 V_s^3 n_s c^4 m^{*2}\sqrt{|\Delta\omega_0|}} \\ &\times \int_0^{2\pi} d\varphi \int_0^{\pi} \sin\theta d\theta |D_{\mathbf{q}}|^2 \exp(-y^2) y_1^2 \\ &\times [L_n^1(y_1^2)]^2 [L_m(y_2^2)]^2 [L_l(y_3^2)]^2 \frac{(n+1)(X_1^i)^2}{(\omega_1-\omega_i)^2+\Gamma^2} \\ &\times \left| \frac{y_2 X_2^S}{\omega_2-\omega_s-i\Gamma}\exp(-i\varphi_2) + \frac{y_3 X_3^S}{\omega_3-\omega_s-i\Gamma}\exp(-i\varphi_3)\right|^2,\end{aligned} \tag{15}$$

Here we replace $N_{\mathbf{q}}$ by the Plank distribution function N_0, y_j can be obtained from g_j if we write the vector $\mathbf{q}$ in the spherical coordinates, $D_{\mathbf{q}}$ depends on the electron-phonon interaction and $\Delta\omega_0 = \omega_1 - \omega_s - \omega_0$.

Let us consider, first of all, the polarization potential scattering (PO phonons). In this case the electron-phonon coupling constant

$$|D_{\mathbf{q}}|^2 = \frac{2\sqrt{2}\pi\hbar^2\alpha_L\omega_0^{3/2}}{\sqrt{m^*}q^2}. \tag{16}$$

Then we can rewrite Equation (15) as

$$\begin{aligned}\frac{d^2\sigma}{d\Omega d\omega_s} &= \frac{\sqrt{2}\pi V\omega_s^2 n_i n_s^3 e^4 N_i(N_s+1)N_0\omega_0^2}{4V_s^3 c^4 m^{*3/2}\varepsilon^2|\Delta\omega_0|^{3/2}} \\ &\int_0^{2\pi} d\varphi \int_0^{\pi} \sin\theta d\theta |D_{\mathbf{q}}|^2 \exp(-y^2) y_1^2 \\ &\times [L_n^1(y_1^2)]^2 [L_m(y_2^2)]^2 [L_l(y_3^2)]^2 \frac{(n+1)(X_1^i)^2}{(\omega_1-\omega_i)^2+\Gamma^2} \\ &\times \left| \frac{y_2 X_2^S}{\omega_2-\omega_s-i\Gamma}\exp(-i\varphi_2) + \frac{y_3 X_3^S}{\omega_3-\omega_s-i\Gamma}\exp(-i\varphi_3)\right|^2,\end{aligned} \tag{17}$$

The cross section depends on the polarization both the input signal and output one. Let us consider the case when the polarization vector of the incident and emitted fields are perpendicular to the magnetic field. In this case the hybrid frequencies are determined by the formulae $\omega_{1,2} = [\sqrt{(\Omega_x+\Omega_y)^2+\omega_c^2} \pm \sqrt{(\Omega_x-\Omega_y)^2+\omega_c^2}]/2$, $\omega_3 = \Omega_z$, and the values of y_j $(j = 1,2,3)$ have the following form

$$\begin{aligned} y_j &= \frac{l_j}{\sqrt{2}}\sqrt{\frac{\omega_0|\Delta\omega_0|}{V_s^2}} \frac{\sin\theta}{\sqrt{\omega_c^2\Omega_x^2 + (\Omega_x^2-\omega_j^2)^2}} \\ &\times \left[(\Omega_x^2-\omega_i^2)^2\sin^2\varphi + \omega_c^2\omega_i^2\cos^2\varphi\right]^{1/2}\end{aligned} \tag{18}$$

Note that in this case $X_j^i = X_j^S$ $(j = 1,2,3)$. Equation (17) clearly shows that the Raman cross section has singularities at the points where $\Delta\omega_0 = 0$. On the Figure 3 it is shown the dependence of the Raman cross section on the magnetic field (here we taken into account the finite phonon relaxation time).

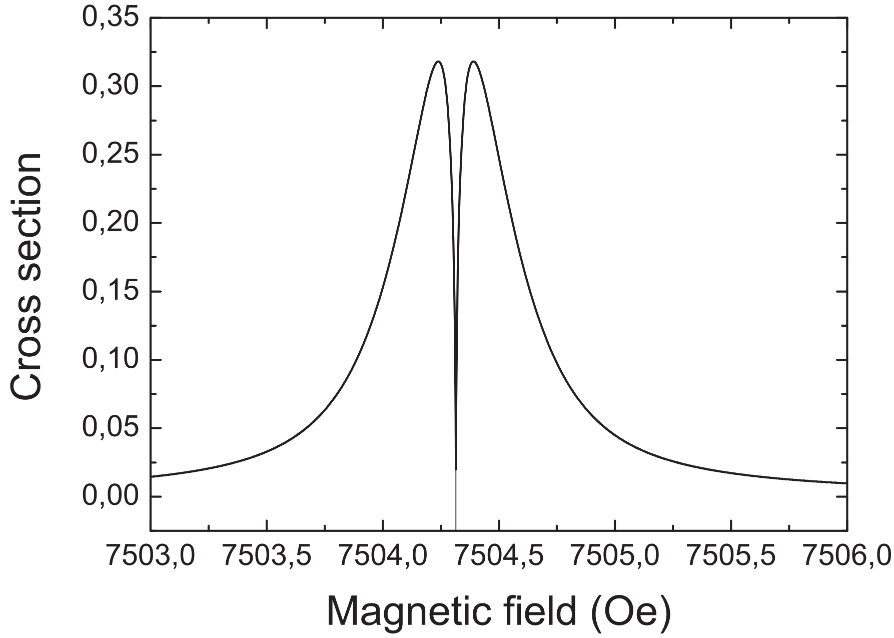

FIG. 4. Differential cross section (in arb. units) as a function of a magnetic field in the case of transition from the ground state and emission of DO-phonons, Other parameters coincide with those of Fig. 3

The different situation takes place in the case of deformation potential scattering (DO phonons). The cross section of deformation potential scattering connected with one of polarization potential scattering by the following estimation

$$\frac{d^2\sigma_{PO}}{d\Omega d\omega_s} = \frac{mV_s^2}{2\hbar|\Delta\omega_0|}\frac{d^2\sigma_{DO}}{d\Omega d\omega_s} \tag{19}$$

It is important to note that in the points where $\Delta\omega_0 = 0$ the differential cross section is equal to zero in contradiction to the case of PO-phonons. In the case of DO-phonons the cross section has the complex doublet structure. The width of the resonance curve is enough small (of order 1 Oe) in this situation. In the most simple case of transitions from the ground state $n = m = l = 0$ the resonance curve consists of two symmetrically positioned sharp peaks to the left and right of the point $\Delta\omega_0 = 0$ (Fig.4). In the case of transitions from the state $n = m = l = 1$ each of the doublet peaks splits up into two ones (Fig.5). If we take into account the finite phonon relaxation time in QDs, the resonance curve doesn't change its form in the difference from the polarization scattering but its minimum displaces in the point where $\Delta\omega + \tau^{-1} = 0$ (here τ is the relaxation time).

In conclusion, we have investigated theoretically the intraband resonance scattering of electromagnetic radiation in the anisotropic quantum dots in the presence of arbitrarily directed magnetic field. We showed that resonance scattering lets us to detect phonon modes in QD using the tunable laser and changing magnetic field. If we ignore optical phonon dispersion, we have a resonance peak corresponding to the emission of optical phonon mode. The interesting doublet structure of peaks arises if one takes into account the dispersion of long-wave optical phonons in the case of deformation scattering. In this case the resonances let us to observe the

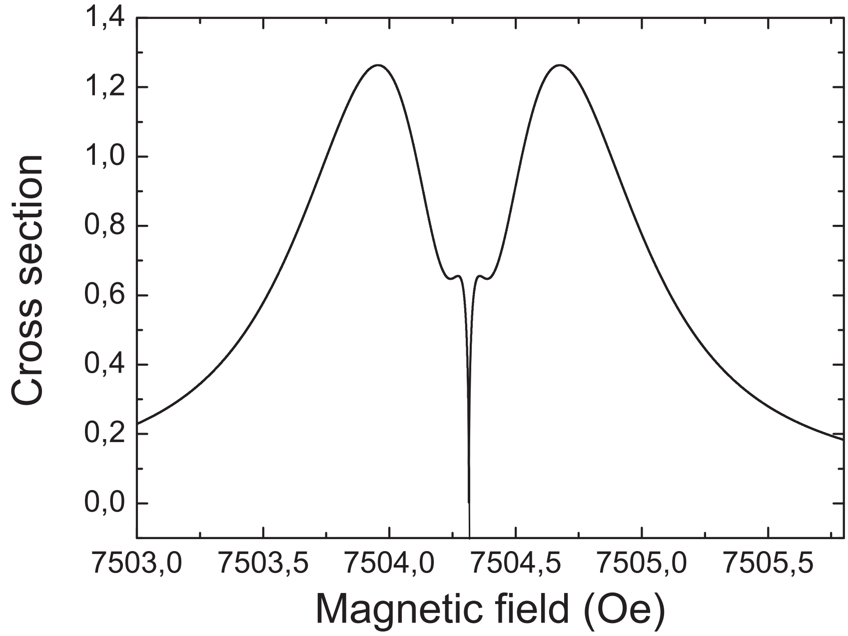

FIG. 5. Differential cross section (in arb. units) as a function of a magnetic field in the case of transition from the state $n = m = l = 1$ and emission of DO-phonons. Other parameters coincide with those of Fig. 3

threshold frequency of optical phonons. In the resonance point the cross section is equal to zero but in a small neighborhood of this point cross section has symmetrically positioned (to the left and right) peaks. The number of peaks depends on the initial quantum state. We hope that our calculations can further stimulate more experimental measurements on the resonance scattering in semiconductor QDs.

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